Resonance-Enhanced Two-Photon Ionization of Ions by Lyman α Radiation in Gaseous Nebulae

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One of the mysteries of nebulae in the vicinity of bright stars is the appearance of bright emission spectral lines of ions, which imply fairly high excitation temperatures. We suggest that an ion formation mechanism, based on resonance-enhanced two-photon ionization (RETPI) by intense H Lyman α radiation (wavelength of 1215 angstroms) trapped inside optically thick nebulae, can produce these spectral lines. The rate of such an ionization process is high enough for rarefied gaseous media where the recombination rate of the ions formed can be 10^{-6} to 10^{-8} per second for an electron density of 10^3 to 10^5 per cubic centimeter in the nebula. Under such conditions, the photo-ions formed may subsequently undergo further RETPI, catalyzed by intense He I and He II radiation, which also gets enhanced in optically thick nebulae that contain enough helium.

Planetary nebulae are known to have no energy sources of their own. The radiation they emit in their bright emission lines is borrowed from the stars embedded within them (*1, 2*). Because of this, photoinduced processes have been known to play a key role in the physics of planetary and gaseous nebulae, as they are responsible for the inflow of energy from the central star(s). Absorption of the short-wavelength stellar blackbody radiation results in the photo-ionization of hydrogen and helium, which repeatedly take part in this process. The ions recombine with electrons, thus providing for a continuous conversion of the absorbed stellar radiation into kinetic energy of the photoelectrons produced and radiative energy of the recombination lines, especially the high-intensity extreme ultraviolet (EUV) lines in the Lyman series of H I and He II and the resonance lines of He I. The radiation of these lines in the optically dense nebula [the physical size (radius) of the nebula is much larger than the free pathway of photons with respect to the resonance scattering events] suffers from diffusion trapping. As a result of the resonance light trapping, there will be a Doppler diffusive broadening of the spectral lines and an increase of their intensity inside the nebula (*1–3*). An increase in intensity of the trapped spectral lines is limited by the radiative cooling of the nebula, caused by the escape of the trapped photons through the wings of the spectral lines (*3*), and the optical excitation of heavier elements by the Bowen mechanism (*4*). The latter is due to an accidental coincidence in wavelength between the trapped spectral lines and particular absorption lines of the heavy elements. Nevertheless, the electron temperature of the nebulae can reach 8×10^3 to 15×10^3 K, and the effective temperature of the highly trapped spectral lines (defined as the temperature of a black body radiating the same spectral intensity at the wavelength of the trapped spectral line) can be at least as high as the electron temperature (*1–3*).

However, it is well known that nebular spectra contain lines from multicharged ions of many elements. For the formation of these ions, it is usually assumed that the central star, which supplies the nebulae with radiative energy, emits EUV radiation with enhanced intensity, which is higher than the equilibrium value corresponding to the temperature of the star photosphere and which can give rise to multicharged ions (*1–3*). Within the framework of the suggested RETPI mechanism, no such assumption is required, as it only requires the summation of two quanta of the Lyman α (Lyα) radiation generated and trapped within the nebula (and not the direct diluted far EUV radiation from the central star, which is too weak). The central star provides only the "pumping" of the nebula by blackbody radiation with $h\nu > 13.6$ eV (h, Planck's constant; ν, frequency), which ensures the photo-ionization of H I and the storage of energy in the form of H II ions and highly trapped recombination radiation Lyα.

The RETPI effect is one version of a two-photon absorption process predicted by Göppert-Mayer (*5*) in 1931 and experimentally verified 30 years later (*6*), following the invention of laser sources of intense coherent light that led to the discovery of numerous nonlinear optical effects (*7, 8*). Almost all of these effects use not only the high intensity of laser light but also its coherence, which makes it possible to accumulate nonlinear effects, specifically to generate optical harmonics (*9*) and other effects in conditions of phase matching. Two-photon absorption, however, requires no light coherence, but it increases substantially in conditions where there exists a quasi-resonance with some intermediate quantum level. This allows such an effect to become manifest in the case of two-photon ionization, despite a low rate on the laboratory scale (*10*), provided that the produced ions accumulate. It is precisely this condition that can be fulfilled in the rarefied medium of a gas nebula in space, where the recombination rate of the ions produced is extremely low (10^{-5} to 10^{-8} s^{-1}) and where an intense EUV radiation exists at the same time in the form of trapped spectral lines of hydrogen and helium.

Figure 1 presents a generalized schematic diagram of the RETPI of an atom (ion) by photons with energies of $h\nu_1$ and $h\nu_2$, whose total energy is $h\nu_1 + h\nu_2 > I_{XN}$ (I_{XN} is the ionization potential of the ion XN, and $N - 1$ is the ionization state). The frequency of the intense radiation, $h\nu_1$, is in quasi-resonance with a frequency $\nu_0 = \nu_{12}$ of an allowed quantum transition 1→2 of the atom (ion), and radiation at the frequency ν_2 (it is possible that $\nu_1 = \nu_2$) is sufficient to ionize the virtually excited X^*N ion. For simplicity, we consider the case where the frequency detuning $\Delta\nu = |\nu_1 - \nu_0| \gg \delta\nu_1$, $\delta\nu_2$, where $\delta\nu_1$ and $\delta\nu_2$ are the spectral widths of the intense radiation at the frequencies ν_1 and ν_2, respectively. The probability of the RETPI

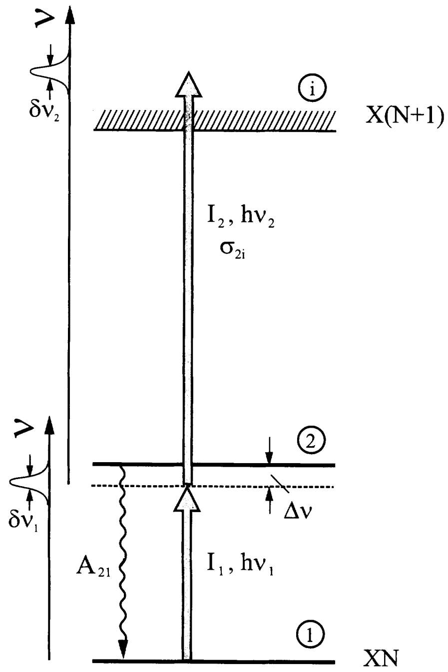

Fig. 1. Generalized schematic diagram of RETPI by intense noncoherent monochromatic radiation at frequencies of ν_1 and ν_2 and spectral widths of $\delta\nu_1$ and $\delta\nu_2$, with a detuning of $\Delta\nu$ from the frequency of exact resonance with the intermediate resonant level *2*.

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process, $W^{(2)}_{1i}$ (s^{-1}), can be derived from (*8*, *10*) to the simple expression

$$W^{(2)}_{1i} = \frac{1}{32\pi^3}\frac{g_2}{g_1}\frac{\lambda_{12}^2}{(\Delta\nu)^2}A_{21}\sigma_{2i}I_1I_2 \qquad (1)$$

where g_i is the degeneracy of the ith level, λ_{12} is the wavelength of the resonance line $1\rightarrow2$ that is close to the wavelength λ_1 of the exciting radiation, A_{21} is the Einstein coefficient of the allowed transition $2\rightarrow1$, σ_{2i} is the photo-ionization cross section of the excited level at the frequency ν_2, and I_1 and I_2 are the radiation intensities of the spectral lines (in photons cm^{-2} s^{-1}), determined by the spectral intensities P_i (photons cm^{-2} s^{-1} Hz^{-1} sr^{-1}) and the spectral widths $\delta\nu_i$ of the trapped lines, $I_i \cong 4\pi P_i \delta\nu_i$. The spectral intensity P_i of the intense lines in the nebula can be described in terms of the effective temperature $T_{\rm eff}$ of the equivalent Planck blackbody radiation as

$$P_i = \frac{2}{\lambda_i^2}\frac{1}{e^{h\nu_i/kT^i_{\rm eff}}-1} \qquad (2)$$

(*e*, base of natural logarithm; *R*, Boltzmann's constant). It must be emphasized that $T_{\rm eff}$ is the effective temperature of the trapped spectral lines inside the nebula, which is much higher than the effective surface temperature of the escaped radiation during the light diffusion process. With this notation, the RETPI probability (Eq. 1) may be represented in the form

$$W_{1i} = \frac{2}{\pi}\frac{g_2}{g_1}\frac{\delta\nu_1\delta\nu_2}{(\Delta\nu)^2}\frac{\sigma_{2i}}{\lambda_2^2} \times A_{21}\frac{1}{(e^{h\nu_1/kT^1_{\rm eff}}-1)}\frac{1}{(e^{h\nu_2/kT^2_{\rm eff}}-1)} \qquad (3)$$

where $\delta\nu_1$ and $\delta\nu_2 \ll \Delta\nu$. Because $h\nu_i$ in the case of Lyα and other, shorter wavelength intense lines (H I, He I, and He II) is much greater than $kT^i_{\rm eff}$, Eq. 3 becomes

$$W_{1i} \cong \frac{2}{\pi}\frac{g_2}{g_1}\frac{\delta\nu_1\delta\nu_2}{(\Delta\nu)^2}\frac{\sigma_{2i}}{\lambda_2^2}A_{21}e^{-\frac{h\nu_1}{kT^1_{\rm eff}}-\frac{h\nu_2}{kT^2_{\rm eff}}} \qquad (4)$$

For the sake of illustration, we consider the RETPI mechanism in connection with the appearance of Si III and Si IV lines in the spectrum of gaseous condensations in the vicinity of η Carinae (*11*), observed with the Hubble Space Telescope. In Fig. 2, we show schematic diagrams of the quantum levels and transitions involved in the RETPI production of Si II, Si III, and Si IV, based on the data presented in (*12*). The mechanism for producing Si II (Fig. 2, left) (by electron collision or photo-ionization by the radiation of the central star or by intense trapped H Lyα radiation) is irrelevant for the RETPI process producing higher silicon ions. H Lyα radiation at λ = 1215 Å is in quasi-resonance with two strong transitions in Si II at λ_1 = 1260 and 1190 Å (*12*), with a frequency detuning of $\Delta\nu$ = 3000 and 1650 cm^{-1}, respectively (Fig. 2, middle).

For a typical case where H Lyα has a temperature of $T_{\rm eff}$ = 15,000 K and a spectral width of $\delta\nu_1 = \delta\nu_2 \cong 300$ cm^{-1} ($\nu_1 = \nu_2$) and the Si II parameters are $A_{21} \cong 10^9$ s^{-1} and $\sigma_{21} \cong 10^{-18}$ cm^2, Eq. 4 yields $W_{1i} \cong 10^{-7}$ s^{-1}. This value is comparable with the recombination rate of the Si II ion at an electron concentration of $n_e \cong 10^5$ cm^{-3}. Thus, the RETPI of Si II in the ground state allows the formation and accumulation of Si III ions.

On the right-hand side of Fig. 2, we show a schematic diagram of RETPI operating on the Si III ion, which has a quasi-resonance (the line at λ = 1206 Å) with H Lyα with a detuning $\Delta\nu$ of 600 cm^{-1}. Photo-ionization of the excited Si III ions requires a shorter wavelength radiation, which may be the strong line of He I at λ_2 = 522 Å. This line is strong in nebular spectra (*1*–*3*) and may have a high temperature $T^2_{\rm eff}$, because of the large optical density and, consequently, the strong optical diffusion trapping. If we use the same parameters for Lyα(H I) as indicated above and assume for λ_2 that $T^2_{\rm eff}$ = 25,000 K and $\delta\nu_2$ = 1000 cm^{-1}, then for $A_{21} \cong 10^9$ s^{-1} and $\sigma_{2i} \cong 10^{-18}$ cm^2, we get, according to Eq. 4, the same estimate of $W_{1i} \cong 10^{-7}$ s^{-1}. Thus, RETPI can proceed consecutively until the formation of such ions that do not have any suitable quasi-resonance with the strong lines of H I, He I, and He II. In the case of silicon ions, this sequence ends with Si IV. Emission lines should be observed for Si III and Si II as a result of recombination of Si IV and Si III, and so on.

To understand the efficiency of the RETPI process, let us compare its rate W_{1i} to that of ionization by electrons. In both cases, there exists the exponential factor exp($-E/kT$), where $E = h(\nu_1 + \nu_2) > IP$ (*IP* is the ionization potential, and *T* is the corresponding radiation or electron temperature). The main difference between ionization through electron collisions and photonic collision-free RETPI is given by the difference in the corresponding nonexponential factors A_e and $A_{\rm ph}$, respectively. According to Eq. 4, in the case of RETPI

$$A_{\rm ph} \cong \frac{2}{\pi}\frac{g_2}{g_1}\frac{\delta\nu_1\delta\nu_2}{(\Delta\nu)^2}\frac{\sigma_{2i}}{\lambda_2^2}A_{21} \qquad (5)$$

which, for the case of Si II and Lyα discussed above, gives $A_{\rm ph} \cong$ (1 to 10) s^{-1}. In the case of electronic excitation, the nonexponential factor is $A_e \cong \langle n_e v_e \sigma_e\rangle$, where n_e is the electron concentration, v_e is the velocity of electrons with energy $E > IP$, and σ_e is the electronic ionization cross section. For $n_e \cong 10^5$ cm^{-3}, we get $A_e \cong 10^{-2}$ to $10^{-3} \ll A_{\rm ph}$. Therefore, the rate of formation of ions by RETPI is higher than that through collisions with energetic electrons and is comparable with the recombination rate of the photo-ions produced. For this reason, the concentration of photo-ions (for example, Si III) produced by RETPI is much higher than that produced by collisions with electrons having kinetic energy $E_{\rm kin} > IP$(Si II), where IP(Si II) is the ionization potential of Si II. Accordingly, the recombination emission lines of Si II become brighter. In addition, the concentration of the ions produced (for example, Si III) becomes sufficient for a subsequent RETPI to Si IV and the appearance of bright emission lines of Si III.

The RETPI is not an exotic, rare process, but it is an effective mechanism for the for-

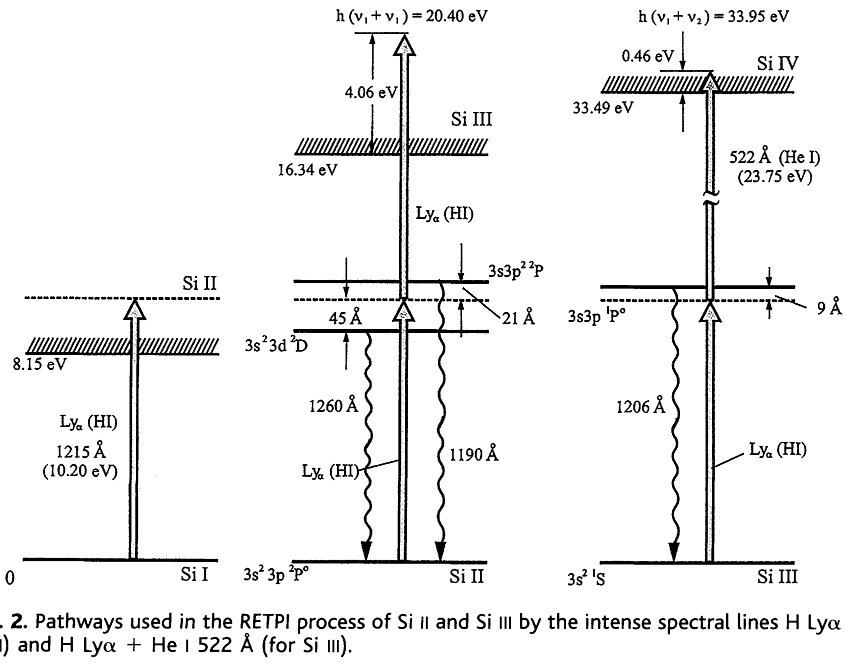

Fig. 2. Pathways used in the RETPI process of Si II and Si III by the intense spectral lines H Lyα (for Si II) and H Lyα + He I 522 Å (for Si III).

mation of multiply charged ions (up to +4) of many elements, such as carbon, nitrogen, oxygen, neon, and argon by strong H I, He I, and He II lines in optically thick nebulae. Moreover, for some ions, the detuning $\Delta\nu$ of frequency ν_1 from the corresponding allowed quantum transition frequency ν_{12} is rather low (<100 cm^{-1}), and the rate of RETPI in the case of such an exact resonance is much higher than for the above considered case of subsequent RETPI in silicon: Si II → Si III → Si IV. We think that the suggested photonic RETPI mechanism is important for understanding the origin of multiply charged ions of many elements inside nebular clouds in the vicinity of bright stars without the requirement of an extremely high (>50,000 K) temperature of the central star's photosphere (*1, 2*).

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Cherenkov Radiation at Speeds Below the Light Threshold: Phonon-Assisted Phase Matching

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Charged particles traveling through matter at speeds larger than the phase velocity of light in the medium emit Cherenkov radiation. Calculations reveal that a given angle of the radiation conical wavefront is associated with two velocities, one above and one below a certain speed threshold. Emission at subluminal but not superluminal speeds is predicted and verified experimentally for relativistic dipoles generated with an optical method based on subpicosecond pulses moving in a nonlinear medium. The dipolar Cherenkov field, in the range of infrared-active phonons, is identical to that of phonon polaritons produced by impulsive laser excitation.

Cherenkov radiation (CR) is extensively used in experiments for counting and identifying relativistic particles (*1, 2*). The effect derives its name from Pavel Cherenkov (*3*), who, following a suggestion by Vavilov (*4*), discovered in 1934 that substances exposed to fast electrons emit coherent light (*5*). The theory of CR was developed by Frank and Tamm in 1937 (*6*). They showed that charges traveling faster than the speed of light in a substance with a frequency-independent refractive index n emit radiation displaying a shock-wave singularity at the surface of a cone defined by $\cos\theta_C = c/nv$, where θ_C is the angle between the direction of motion and that of the electromagnetic wavefront (the angle subtended by the cone is $\alpha_C = \pi/2 - \theta_C$), c is the speed of light in vacuum, and v is the speed of the particle.

We concern ourselves with CR in an isotropic material whose frequency (Ω)–dependent dielectric constant $\varepsilon = n^2$ in the vicinity of a resonance at frequency Ω_0 can be approximated by the nondissipative Lorentz form

$$\varepsilon(\Omega) = \varepsilon_\infty + \frac{\varepsilon_0 - \varepsilon_\infty}{1 - (\Omega/\Omega_0)^2} \qquad (1)$$

ε_0 is the dielectric constant at $\Omega = 0$, and ε_∞ accounts for the contribution of higher lying excitations. Although analytical expressions for the Cherenkov field of a generic dielectric have been available for a long time (*7*), its properties in the case of a strongly dispersive substance have not been studied in any detail until very recently (*8*). Here, we report on both experimental and theoretical work revealing important qualitative differences for speeds below and above the speed of light at $\Omega = 0$, $c_0 = c/n(0)$ ($v > c_0$ and $v < c_0$ are referred to in the following as the superluminal and subluminal regimes) (*9*). We use subpicosecond optical pulses to generate an optical analog of a relativistic dipole (*10, 11*). Measurements of the Cherenkov electric field **E** for a planar distribution of such dipoles show that the field is nonzero for $v < c_0$ but vanishes for $v > c_0$. We also revisit the standard problem of CR emitted by a point charge and find that, regardless of v, the field pattern is always outlined by a cone. However, the v dependence of θ_C exhibits a minimum $\theta_C = 0$ at $v = c_0$, contradicting the widely held belief that there is a one-to-one correspondence between v and the Cherenkov angle (*1, 2*). Finally, we reinterpret experiments on the ultrafast generation of polaritons (*12–17*) in terms of CR.

A representative set of our calculations is reproduced in Fig. 1. $\varepsilon(\Omega)$ is given by Eq. 1 and the radiation source is moving along the z axis. The magnitude of the field E for a point particle is shown in Fig. 1A (*18*). Because of the symmetry of the problem, the electromagnetic field depends on $z - vt$ (t is the time) and $\rho = (x^2 + y^2)^{1/2}$, but not on the azimuth. The diagrams in Fig. 1B depict E for a collection of electric dipoles uniformly distributed in a plane. Both the dipole moment and the lane are perpendicular to the velocity. The latter case is directly relevant to our experiment.

Consider first the point source. The features of interest are (i) the observation that θ_C is well defined at all speeds (*19*), (ii) the fast oscillations at $v > c_0$ with a frequency that increases with decreasing $\alpha = \tan^{-1}\rho/(vt - z)$, and (iii) the subluminal beats. More importantly, our results indicate that a given Cherenkov angle is connected not with one, but with two particle speeds (one below and one above the light threshold).

For a point charge, the electromagnetic potentials can be expressed in terms of integrals of the form (*7*)

$$A(t - z/v, \rho, v) = \int \Phi(\Omega) e^{-i\Omega(t-z/v)} H_0^{(1)}(s\rho) d\Omega \qquad (2)$$

$H_0^{(1)}$ is a Hankel function, Φ is a slowly varying function of Ω, and $s^2(\Omega) = \Omega^2(\varepsilon v^2/c^2 - 1)/v^2$. At large distances, we can restrict the integration to the radiative range $\Omega_C \leq \Omega \leq \Omega_0$ where $s^2 \geq 0$; Ω_C is either the frequency at which $v^2 = c^2/\varepsilon$ ($v < c_0$) or $\Omega_C = 0$ ($v > c_0$). Using the asymptotic expression of $H_0^{(1)}$,

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